

Parity doublets in the baryon spectrum.

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Physics of the low-lying and high-lying hadrons in the light flavor sector is reviewed. While the low-lying hadrons are strongly affected by both $U(1)_A$ and spontaneous $SU(2)_L \times SU(2)_R$ breakings, in the high-lying hadrons these symmetries are restored. A manifestation is a persistence of the chiral multiplet structure in both baryon and meson spectra. A fundamental origin of this phenomenon is that effects of quantum fluctuations of both quark and gluon fields must vanish at large n or J and a semiclassical description becomes adequate. A relation between the chiral symmetry restoration and the string picture of excited hadrons is discussed.

1. Introduction

If one neglects the tiny masses of the u and d quarks, which are much smaller than Λ_{QCD} , then the QCD Lagrangian exhibits the

$$U(2)_L \times U(2)_R = SU(2)_L \times SU(2)_R \times U(1)_V \times U(1)_A \quad (1)$$

symmetry. This is because the quark-gluon interaction Lagrangian in the chiral limit does not mix the left- and right-handed components of quarks and hence the total QCD Lagrangian for the two-flavor QCD can be split into left-handed and right-handed parts which do not communicate to each other. We know that the $U(1)_A$ symmetry of the classical QCD Lagrangian is absent at the quantum level because of the $U(1)_A$ anomaly, which is an effect of quantum fluctuations [1]. We also know that the chiral $SU(2)_L \times SU(2)_R$ symmetry is spontaneously (dynamically) broken in the QCD vacuum [2]. That this is so is directly evidenced by the nonzero value of the quark condensate, $\langle \bar{q}q \rangle = \langle \bar{q}_L q_R + \bar{q}_R q_L \rangle \simeq -(240 \pm 10 \text{ MeV})^3$, which represents an order parameter for spontaneous chiral symmetry breaking. This quark condensate shows that in the QCD vacuum the left-handed quarks are correlated with the right-handed antiquarks (and vice versa) and hence the QCD vacuum breaks the chiral symmetry. This spontaneous (dynamical) breaking of chiral symmetry is a pure quantum effect based upon quantum fluctuations. As a consequence we do not observe any chiral or $U(1)_A$ multiplets low in the hadron spectrum.

The upper part of both baryon [3, 4] and meson [5] spectra almost systematically exhibits multiplets of the chiral and $U(1)_A$ groups (for a pedagogical overview see [6]), though a careful experimental exploration of high-lying spectra must be done for a final conclusion. This phenomenon is referred to as effective chiral symmetry restoration or chiral symmetry restoration of the second kind. This is illustrated in Fig. 1, where the excitation spectrum of the nucleon from the PDG compilation as well as the excitation spectrum of π and f_0 (with the $\bar{n}n = \frac{\bar{u}u + \bar{d}d}{\sqrt{2}}$ content) mesons [5] are shown. Starting from the 1.7 GeV region the nucleon (and delta) spectra show obvious signs of parity doubling. There are a couple of examples where chiral partners

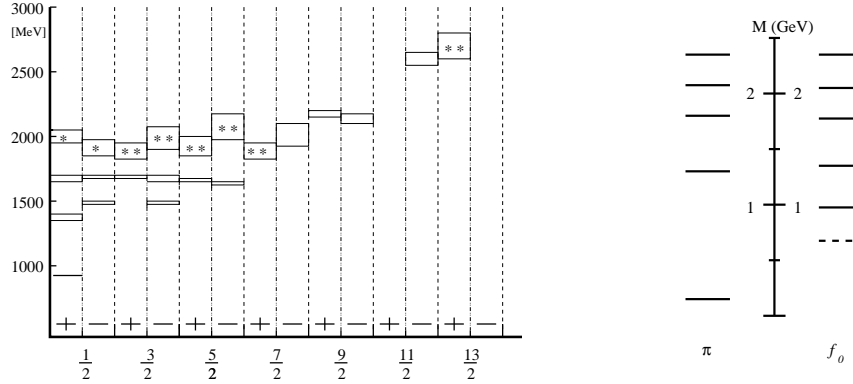


Figure 1. Left panel: excitation spectrum of the nucleon (those resonances which are not yet established are marked by two or one stars according to the PDG classification). Right panel: pion and $n\bar{n} f_0$ spectra.

of highly excited states have not yet been seen. Their experimental discovery would be an important task. Similarly, in the chirally restored regime π and $n\bar{n} f_0$ states must be systematically degenerate.

2. A few words about chiral symmetry breaking and low-lying hadrons

A key to the understanding of the low-lying hadrons is the spontaneous breaking of chiral symmetry (SBCS). Any interquark interaction in QCD mediated by the intermediate gluon field, in the local approximation and once antisymmetrization of the quark fields has been performed, contains as a part a chiral-invariant 4-fermion interaction [2], $(\bar{\psi}\psi)^2 + (\bar{\psi}i\gamma_5\vec{\tau}\psi)^2$. The first term represents a Lorentz-scalar interaction. This interaction is an attraction between the left-handed quarks and the right-handed antiquarks and vice versa. When it is treated nonperturbatively in the mean-field approximation, which is well justified in the vacuum state, it leads to the condensation of the chiral pairs in the vacuum state, $\langle 0|\bar{\psi}\psi|0\rangle = \langle 0|\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L|0\rangle \neq 0$. Hence it breaks chiral symmetry, which is a nonperturbative phenomenon. This attractive interaction between bare quarks can be absorbed into a mass of a quasiparticle. Each quasiparticle is a coherent superposition of bare quarks and antiquarks. Bare particles have both well-defined helicity and chirality, while quasiparticles have only definite helicity and contain a mixture of bare quarks and antiquarks with opposite chirality. These quasiparticles with dynamical mass can be associated with the constituent quarks. An important feature is that this dynamical mass appears only at low momenta, below the ultraviolet cutoff Λ in the NJL model, i.e. where the low-momentum attractive interaction between quarks is operative. All quarks with momenta higher than Λ remain undressed. In reality, of course, this step-function behaviour of the dynamical mass should be substituted by some smooth function.

Once the chiral symmetry is spontaneously broken, then there must appear collective massless Goldstone excitations. Microscopically their zero mass is provided by the term $(\bar{\psi}i\gamma_5\vec{\tau}\psi)^2$. This term represents an attraction between the constituent quark and the antiquark with the pion quantum numbers. Without this term the pion would have a mass of $2M$. When this term is

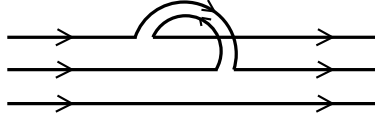


Figure 2. Pion-exchange between valence quarks in the low-lying baryons.

nonperturbatively and relativistically iterated the attraction between the constituent quarks in the pion exactly compensates the $2M$ energy and the pion becomes massless. This happens because of the underlying chiral symmetry since it is this symmetry which dictates that the strengths of the interactions represented by the first and by the second terms in the NJL Hamiltonian are equal. So the pion is a relativistic bound state of two quasiparticles. It contains $\bar{Q}Q$, $\bar{Q}Q\bar{Q}Q$, ... Fock components. The pion (as any Goldstone boson) is a highly collective excitation in terms of the original (bare) quarks and antiquarks q and \bar{q} because the quasiparticles Q and \bar{Q} themselves are coherent collective excitations of bare quarks.

Now we will go to the low-lying baryons. A basic ingredient of the chiral quark picture of Manohar and Georgi [7] is that the constituent quarks inside the nucleon are strongly coupled to the pion field and this coupling is regulated by the Goldberger-Treiman relation. Why this must be so can be seen directly from the Nambu and Jona-Lasinio mechanism of chiral symmetry breaking. Then in the low-momentum regime (which is responsible for masses) the *low-lying* baryons in the light flavor sector can be approximated as systems of three confined constituent quarks with the residual interaction mediated by the Goldstone boson field [8]. Such a model was designed to solve a problem of the low-lying baryon spectroscopy. Microscopically this residual interaction appears from the t-channel iterations of those gluonic interactions in QCD which are responsible for chiral symmetry breaking [9], see Fig. 2. An essential feature of this residual interaction is that it is a flavor- and spin-exchange interaction and contributes to the baryon energy at the order N_c while the contribution of the color-magnetic interaction appears only at $1/N_c$. This specific form of the residual interaction between valence constituent quarks in baryons allows us not only to generate octet-decuplet splittings but what is important to solve at the same time the long-standing puzzle of the ordering of the lowest excitations of positive and negative parity in the u, d, s sector. This physics is a subject of intensive lattice studies and recent results [10, 11, 12, 13] do show that the correct ordering can be achieved only close to the chiral limit and hence is related to spontaneous breaking of chiral symmetry. The results [12] also evidence a node in the wave function of the radial excitation of the nucleon (Roper resonance) which is consistent with the 3Q leading Fock component of this state. It is also important to realize for the following discussion, that all these effects of the pion cloud are effects of *quantum fluctuations of the quark fields*, which is well seen from Fig. 2.

This physics and effective degrees of freedom, which are based on the spontaneous breaking of chiral symmetry, are relevant, however, only to the low-lying hadrons. In the high-lying hadrons chiral symmetry is restored.

3. Chiral symmetry restoration in excited hadrons by definition

By definition chiral symmetry restoration means the following. In QCD hadrons with quantum numbers α are created when one applies the local interpolating field (current) J_α with such quantum numbers on the vacuum $|0\rangle$. Then all hadrons that are created by the given interpolator appear as intermediate states in the two-point correlator,

$$\Pi = i \int d^4x e^{iqx} \langle 0 | T \{ J_\alpha(x) J_\alpha^\dagger(0) \} | 0 \rangle, \quad (2)$$

where all possible Lorentz and Dirac indices (specific for a given interpolating field) have been omitted. Consider two local interpolating fields $J_1(x)$ and $J_2(x)$ which are connected by a chiral transformation (or by a $U(1)_A$ transformation), $J_1(x) = U J_2(x) U^\dagger$. Then, if the vacuum was invariant under the chiral group, $U|0\rangle = |0\rangle$, it follows from (2) that the spectra created by the operators $J_1(x)$ and $J_2(x)$ would be identical. We know that in QCD one finds $U|0\rangle \neq |0\rangle$. As a consequence the spectra of the two operators must be in general different. However, it happens that the noninvariance of the vacuum becomes unimportant (irrelevant) high in the spectrum. Then the spectra of both operators become close at large masses and asymptotically identical. This means that chiral symmetry is effectively restored. We stress that this effective chiral symmetry restoration does not mean that chiral symmetry breaking in the vacuum disappears, but that the role of the quark condensates that break chiral symmetry in the vacuum becomes progressively less important high in the spectrum. One could say, that the valence quarks in high-lying hadrons *decouple* from the QCD vacuum. In order to avoid a confusion with the chiral symmetry restoration in the vacuum state at high temperature or density one also refers this phenomenon as *chiral symmetry restoration of the second kind*.

4. The quark-hadron duality and chiral symmetry restoration

A question arises to which extent the chiral symmetry restoration of the second kind can be theoretically predicted in QCD. There is a heuristic argument that supports this idea [4]. The argument is based on the well controlled behaviour of the two-point function (2) at the large space-like momenta $Q^2 = -q^2$, where the operator product expansion (OPE) is valid and where all nonperturbative effects can be absorbed into condensates of different dimensions [14]. The key point is that all nonperturbative effects of the spontaneous breaking of chiral symmetry at large Q^2 are absorbed into the quark condensate $\langle \bar{q}q \rangle$ and other quark condensates of higher dimension. However, the contribution of these condensates to the correlation function is regulated by the Wilson coefficients. The latter ones are proportional to $(1/Q^2)^n$, where the index n is determined by the quantum numbers of the current J and by the dimension of the given quark condensate. Hence, at large enough Q^2 the two-point correlator becomes approximately chirally symmetric. At these high Q^2 a matching with the perturbative QCD (where no SBCS occurs) can be done. In other words, though chiral symmetry is broken in the vacuum and all chiral noninvariant condensates are not zero, their influence on the correlator at asymptotically high Q^2 vanishes. This is in contrast to the situation of low values of Q^2 , where the role of chiral symmetry breaking in the vacuum is crucial. Hence, at $Q^2 \rightarrow \infty$ one has

$$\Pi_{J_1}(Q^2) - \Pi_{J_2}(Q^2) \sim \frac{1}{Q^n}, \quad n > 0. \quad (3)$$

Now we can use the causality of the local field theory and hence the analyticity of the two-point function. Then we can invoke into analysis a dispersion relation,

$$\Pi_J(Q^2) = \int ds \frac{\rho_J(s)}{Q^2 + s - i\epsilon}, \quad \rho_J(s) \equiv \frac{1}{\pi} \text{Im} (\Pi_J(s)). \quad (4)$$

Since the large Q^2 asymptotics of the correlator is given by the leading term of the perturbation theory, then the asymptotics of $\rho(s)$ at $s \rightarrow \infty$ must also be given by the same term of the perturbation theory if the spectral density approaches a constant value (if it oscillates, then it must oscillate around the perturbation theory value). Hence both spectral densities $\rho_{J_1}(s)$ and $\rho_{J_2}(s)$ at $s \rightarrow \infty$ must approach the same value and the spectral function becomes chirally symmetric. This theoretical expectation, that the high s asymptotics of the spectral function is well described by the leading term of the perturbation theory has been tested e.g. in the process $e^+e^- \rightarrow \text{hadrons}$, where the interpolator is given by the usual electromagnetic vector current. Similarly, the vector and the axial vector spectral densities must coincide in the chiral symmetry restored regime. They have been measured in the τ decay by the ALEPH and OPAL collaborations at CERN [15, 16]. It is well seen from the results that while the difference between both spectral densities is large at the masses of the $\rho(770)$ and $a_1(1260)$, it becomes strongly reduced towards $m = \sqrt{s} \sim 1.7$ GeV.

The question arises then what is the functional behaviour that determines approaching the chiral-invariant regime at large s ? Naively one would expect that the operator product expansion of the two-point correlator, which is valid in the deep Euclidean domain, could help us. This is not so, however, for two reasons. First of all, we know phenomenologically only the lowest dimension quark condensate. Even though this condensate dominates as a chiral symmetry breaking measure at the very large space-like Q^2 , at smaller Q^2 the higher dimensional condensates, which are suppressed by inverse powers of Q^2 , are also important. These condensates are not known, unfortunately. But even if we knew all quark condensates up to a rather high dimension, it would not help us. This is because the OPE is only an asymptotic expansion [18]. While such kind of expansion is very useful in the space-like region, it does not define any analytical solution which could be continued to the time-like region at finite s . While convergence of the OPE can be improved by means of the Borel transform and it makes it useful for SVZ sum rules for the low-lying hadrons, this cannot be done for the higher states. So in order to estimate chiral symmetry restoration effects one indeed needs a microscopic theory that would incorporate *at the same time* chiral symmetry breaking and confinement.

5. Chiral and $U(1)_A$ restorations as a manifestation of the semiclassical regime

While the argument above on the asymptotic symmetry properties of spectral functions is rather robust (it is based actually only on the asymptotic freedom of QCD at large space-like momenta and on the analyticity of the two-point correlator), *a-priori* it is not clear whether it can be applied to bound state systems, which the hadrons are. Indeed, it can happen that the asymptotic symmetry restoration applies only to that part of the spectrum, which is above the resonance region (i.e. where the current creates jets but not isolated hadrons). So the question arises whether it is possible to prove (or at least justify) the symmetry restoration in highly excited *isolated* hadrons. It is shown in ref. [17] that both chiral and $U(1)_A$ restorations in highly excited isolated hadrons can be anticipated as a direct consequence of the semiclassical regime in the highly excited hadrons, indeed.

At large n (radial quantum number) or at large angular momentum J we know that in quantum systems the *semiclassical* approximation (WKB) *must* work. Physically this approximation applies in these cases because the de Broglie wavelength of particles in the system is small in comparison with the scale that characterizes the given problem. In such a system as a hadron the scale is given by the hadron size while the wavelength of valence quarks is given by their momenta. Once we go high in the spectrum the size of hadrons increases as well as the typical momentum of valence quarks. This is why a highly excited hadron can be described semiclassically in terms of the underlying quark degrees of freedom.

A physical content of the semiclassical approximation is most transparently given by the path integral. The contribution of the given path to the path integral, $\sim e^{iS(q)/\hbar}$, is regulated by the action $S(q)$ along the path $q(x, t)$. The semiclassical approximation applies when the action in the system $S \gg \hbar$. In this case the whole amplitude (path integral) is dominated by the classical path q_{cl} (stationary point) and those paths that are infinitesimally close to the classical path. All other paths that differ from the classical one by an appreciable amount do not contribute. These latter paths would represent the quantum fluctuation effects. In other words, in the semiclassical case the quantum fluctuations effects are strongly suppressed and vanish asymptotically. The classical path is a tree-level contribution to the path integral and keeps all symmetries of the classical Lagrangian. Its contribution is of the order S/\hbar . The quantum fluctuations of the fields contribute at the orders $(\hbar/S)^0$ (one loop), $(\hbar/S)^1$ (two loops), etc.

The $U(1)_A$ as well as spontaneous $SU(2)_L \times SU(2)_R$ breakings result from quantum fluctuations of the quark fields. However, in a quantum system with large enough n or J the quantum fluctuations contributions must be relatively suppressed and vanish asymptotically. Then it follows that in such systems both the chiral and $U(1)_A$ symmetries must be restored. Hence at large hadron masses (i.e. with either large n or large J) we must observe the symmetries of the classical QCD Lagrangian. This is precisely what we see phenomenologically. In the nucleon spectrum the doubling appears either at large n excitations of baryons with the given small spin or in resonances of large spin. Similar features persist in the delta spectrum. In the meson spectrum the doubling is obvious for large n excitations of small spin mesons and there are signs of doubling of large spin mesons (the data are, however, sparse). It would be certainly interesting and important to observe systematically multiplets of parity-chiral and parity- $U(1)_A$ groups (or, sometimes, when the chiral and $U(1)_A$ transformations connect *different* hadrons, the multiplets of the $U(2)_L \times U(2)_R$ group [5]).

While the argument above is robust, theoretically it is not clear *a-priori* whether isolated hadrons still exist at excitation energies where a semiclassical regime is achieved. Hence it is conceptually important to demonstrate that in QCD hadrons still exist, while the dynamics inside such hadrons already is semiclassical. We do not know how to prove it for $N_c = 3$. However, the large N_c limit of QCD [19], while keeping all basic properties of QCD like asymptotic freedom and chiral symmetry, allows for a significant simplification. In this limit it is known that all mesons represent narrow states, i.e. they are stable against strong decays. At the same time the spectrum of mesons is infinite (see, e.g., [20]). The latter is necessary to match the two-point function in the perturbation theory regime (which contains logarithm) at large space-like momenta with the discrete spectral sum in the dispersion integral. Then one can always excite a meson of any arbitrary large energy, which is of any arbitrary large size. In such a meson $S \gg \hbar$. Hence a description of this meson necessarily must be semiclassical. Then the equation of motion in such a meson must be according to some yet unknown solution

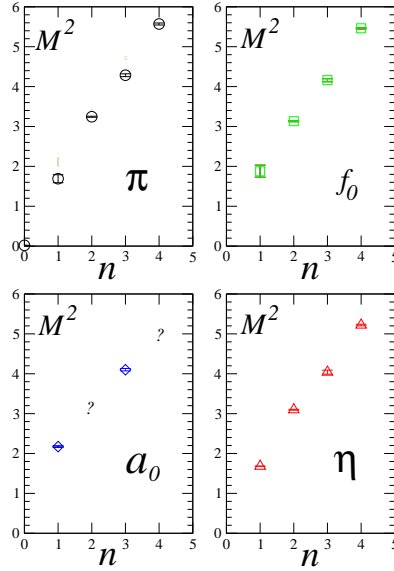


Figure 3. Radial Regge trajectories for the four successive high-lying $J = 0$ mesons.

of the classical QCD Lagrangian for a colorless meson. Hence it must exhibit chiral and $U(1)_A$ symmetries.

Actually we do not need the exact $N_c = \infty$ limit for this statement. It can be formulated in the following way. For any large $S \gg \hbar$ there always exist such N_c that the isolated meson with such an action does exist and can be described semiclassically. From the empirical fact that we observe multiplets of chiral and $U(1)_A$ groups high in the hadron spectrum it follows that $N_c = 3$ is large enough for this purpose.¹

¹That the quantum fluctuations effects vanish in the quantum bound state systems at large n or J is well known e.g. from the Lamb shift. The Lamb shift is a result of the radiative corrections (which represent effects of quantum fluctuations of electron and electromagnetic fields) and vanishes as $1/n^3$, and also very fast with increasing J . As a consequence high in the hydrogen spectrum the symmetry of the classical Coulomb potential gets restored. The other well-known example which clearly illustrates the point is the 't Hooft model [21] (QCD in 1+1 dimensions). In this model in the regime $N_c \rightarrow \infty$, $m_q \rightarrow 0$, $m_q \gg g \sim 1/\sqrt{N_c}$, the spectrum of the high-lying states is known exactly, $M_n^2 \sim n$. The chiral symmetry of the Lagrangian is broken (with no contradiction with the Coleman theorem since in this specific regime everything is determined by $N_c = \infty$, for any large but finite N_c the chiral symmetry is not broken in agreement with the Coleman theorem), which is reflected by the fact that the positive and negative parity states are not degenerate and alternate in the spectrum. However, the mass difference between the neighbouring positive and negative parity states is $m_+ - m_- \sim 1/\sqrt{n}$ and vanishes high in the spectrum since the effect of quantum fluctuations dies out high in the spectrum.

6. Chiral multiplets of excited mesons

A detailed classification of the chiral and $U(1)_A$ multiplets for excited mesons, based on the recent experimental data obtained from proton-antiproton annihilation at LEAR [22, 23], is given in ref. [5]. Here we limit ourselves only to a few typical and spectacular examples.

Consider first the mesons of spin $J = 0$, which are the π , f_0 , a_0 and η mesons with u, d quark content only. If one looks at the upper part of the meson spectrum, then one notices that the four successive excited π mesons and the corresponding $\bar{n}n$ f_0 mesons form approximate chiral pairs and fill out $(1/2, 1/2)$ representations of the chiral group. This is well seen from Fig. 1. This pattern is a clear manifestation from chiral symmetry restoration.

A similar behaviour is observed from a comparison of the a_0 and η masses. However, there are two missing a_0 mesons which must be discovered in order to complete all chiral multiplets. There is little doubt that these missing a_0 mesons do exist. If one puts the four high-lying π , $\bar{n}n$ f_0 , a_0 and $\bar{n}n$ η mesons on the *radial* Regge trajectories, see Fig. 3, one clearly notices that the two missing a_0 mesons lie on the linear trajectory with the same slope as all other mesons. If one reconstructs these missing a_0 mesons according to this slope, then a pattern of the $a_0 - \eta$ chiral partners appears, similar to the one for the π and f_0 mesons.

For the $J \geq 1$ mesons the classification is a bit more complicated. Below we show the chiral patterns for the $J = 2$ mesons, where the data set seems to be complete (masses are in MeV).

(0,0)	
$\omega_2(0, 2^{--})$	$f_2(0, 2^{++})$
1975 ± 20	1934 ± 20
2195 ± 30	2240 ± 15
(1/2,1/2)	
$\pi_2(1, 2^{-+})$	$f_2(0, 2^{++})$
2005 ± 15	2001 ± 10
2245 ± 60	2293 ± 13
(1/2,1/2)	
$a_2(1, 2^{++})$	$\eta_2(0, 2^{-+})$
2030 ± 20	$2030 \pm ?$
2255 ± 20	2267 ± 14
(0,1)+(1,0)	
$a_2(1, 2^{++})$	$\rho_2(1, 2^{--})$
1950^{+30}_{-70}	1940 ± 40
2175 ± 40	2225 ± 35

We see systematic patterns of chiral symmetry restoration. In particular, the amount of $f_2(0, 2^{++})$ mesons coincides with the combined amount of $\omega_2(0, 2^{--})$ and $\pi_2(1, 2^{-+})$ states. Similarly, the number of $a_2(1, 2^{++})$ states is the same as the number of $\eta_2(0, 2^{-+})$ and $\rho_2(1, 2^{--})$ together. All chiral multiplets are complete. While the masses of some of the states can and will be corrected in future experiments, if new states might be discovered in this energy region in other types of experiments, they should be either $\bar{s}s$ states or glueballs.

It is important to see whether there are also signatures of the $U(1)_A$ restoration. Those interpolators (states) that are members of the $(0, 0)$ and $(0, 1) + (1, 0)$ representations of $SU(2)_L \times SU(2)_R$ are invariant with respect to $U(1)_A$. However, interpolators (states) from the *distinct* $(1/2, 1/2)$ representations which have the same isospin but opposite parity transform into each other under $U(1)_A$. If the corresponding states are systematically (approximately) degenerate, then it is a signal that $U(1)_A$ is restored. In what follows we show that it is indeed the case.

$f_2(0, 2^{++})$	$\eta_2(0, 2^{-+})$
2001 ± 10	$2030 \pm ?$
2293 ± 13	2267 ± 14
$\pi_2(1, 2^{-+})$	$a_2(1, 2^{++})$
2005 ± 15	2030 ± 20
2245 ± 60	2255 ± 20

We see clear approximate doublets of $U(1)_A$ restoration. Hence two distinct $(1/2, 1/2)$ multiplets of $SU(2)_L \times SU(2)_R$ can be combined into one multiplet of $U(2)_L \times U(2)_R$. So we conclude that the whole chiral symmetry of the QCD Lagrangian $U(2)_L \times U(2)_R$ gets approximately restored high in the hadron spectrum.

It is useful to quantify the effect of chiral symmetry breaking (restoration). An obvious parameter that characterizes effects of chiral symmetry breaking is a relative mass splitting within the chiral multiplet. Let us define the *chiral asymmetry* as

$$\chi = \frac{|M_1 - M_2|}{(M_1 + M_2)}, \quad (5)$$

where M_1 and M_2 are masses of particles within the same multiplet. This parameter gives a quantitative measure of chiral symmetry breaking at the leading (linear) order and has the interpretation of the part of the hadron mass due to chiral symmetry breaking. For the low-lying states the chiral asymmetry is typically 0.2 - 0.6 which can be seen e.g. from a comparison of the $\rho(770)$ and $a_1(1260)$ or the $\rho(770)$ and $h_1(1170)$ masses. If the chiral asymmetry is large as above, then it makes no sense to assign a given hadron to the chiral multiplet since its wave function is a strong mixture of different representations and we have to expect also large *nonlinear* symmetry breaking effects. However, at meson masses of about 2 GeV the chiral asymmetry is typically within 0.01 and in this case the hadrons can be believed to be members of multiplets with a tiny admixture of other representations.

7. Chiral multiplets of excited baryons

Now we will consider chiral multiplets of excited baryons [4]. The only possible representations of the $SU(2)_L \times SU(2)_R$ group that are invariant under parity are $(1/2, 0) \oplus (0, 1/2)$, $(1/2, 1) \oplus (1, 1/2)$ and $(3/2, 0) \oplus (0, 3/2)$. Since chiral symmetry and parity do not constrain the possible spins of the states these multiplets can correspond to states of any fixed spin.

A phenomenological consequence of the effective restoration of chiral symmetry high in N and Δ spectra is that the baryon states will fill out the irreducible representations of the parity-chiral group. If $(1/2, 0) \oplus (0, 1/2)$ and $(3/2, 0) \oplus (0, 3/2)$ multiplets were realized in nature, then

the spectra of highly excited nucleons and deltas would consist of parity doublets. However, the energy of the parity doublet with given spin in the nucleon spectrum *a-priori* would not be degenerate with the the doublet with the same spin in the delta spectrum; these doublets would belong to different representations, *i.e.* to distinct multiplets and their energies are not related. On the other hand, if $(1/2, 1) \oplus (1, 1/2)$ were realized, then the high-lying states in the N and Δ spectra would have a N parity doublet and a Δ parity doublet with the same spin and which are degenerate in mass. In either case the high-lying spectrum must systematically consist of parity doublets.

If one looks carefully at the nucleon spectrum, see Fig. 1, and the delta spectrum one notices that the systematic parity doubling in the nucleon spectrum appears at masses of 1.7 GeV and above, while the parity doublets in the delta spectrum insist at masses of 1.9 GeV and higher. This means that the parity doubling in both cases is seen at approximately the same excitation energy with respect to the corresponding ground state. This is because not an absolute value of the energy is important in order to approach the chiral symmetry restoration regime but rather n or J . In both nucleon and delta spectra parity doubling persists at the same values of n and J . This fact implies that at least those nucleon doublets that are seen at ~ 1.7 GeV belong to the $(1/2, 0) \oplus (0, 1/2)$ representation. If approximate mass degeneracy between some of the N and Δ doublets at $M \geq 1.9$ GeV is accidental, then the baryons in this mass region are organized according to $(1/2, 0) \oplus (0, 1/2)$ for N and $(3/2, 0) \oplus (0, 3/2)$ for Δ parity-chiral doublets. If not, then some of the high-lying doublets form $(1/2, 1) \oplus (1, 1/2)$ multiplets. It can also be possible that in the narrow energy interval more than one parity doublet in the nucleon and delta spectra is found for a given spin. This would then mean that different doublets belong to different parity-chiral multiplets. Systematic experimental exploration of the high-lying states is required in order to assign unambiguously baryons to the multiplets.

8. Chiral symmetry restoration and the string (flux tube) picture

Before discussing a model for highly excited hadrons that is compatible with the chiral symmetry restoration and parity doubling it is useful to answer the question whether the nonrelativistic or relativized potential models based on the $^{2S+1}L_J$ description like the traditional constituent quark model can explain it. This question can be definitely answered "no"; for a transparent explanation see the review [6].

We have already discussed that in highly excited hadrons the valence quark motion has to be described semiclassically and at the same time their chirality (helicity) must be fixed. Also the gluonic field should be described semiclassically. All this gives an increasing support for a string picture [24, 25] of highly excited hadrons. Indeed, if one assumes that the quarks at the ends of the string have definite chirality, see Fig. 4, then all hadrons will appear necessarily in chiral and $U(1)_A$ multiplets [26]. The latter picture is very natural and is well compatible with the Nambu string picture. The ends of the string in the Nambu model move with the velocity of light. Then, (it is an extension of the Nambu model) the quarks at the ends of the string must have definite chirality. In this way one is able to explain at the same time both Regge trajectories and parity doubling.

There is one additional and important byproduct of this picture [26]. The spin-orbit operator $\vec{\sigma} \cdot \vec{L}$ does not commute with the helicity operator $\vec{\sigma} \cdot \vec{\nabla}$. Hence the spin-orbit interaction of

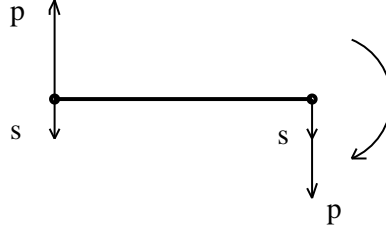


Figure 4. Rotating string with the right and the left quarks at the ends.

quarks with the fixed chirality or helicity is absent. In particular, this is also true for the spin-orbit force due to the Thomas precession

$$U_T = -\vec{\sigma} \cdot \vec{\omega}_T \sim \vec{\sigma} \cdot [\vec{v}, \vec{a}] \sim \vec{v} \cdot [\vec{v}, \vec{a}] = 0, \quad (6)$$

where U_T is the energy of the interaction and $\vec{\omega}_T$, \vec{v} and \vec{a} are the angular frequency of Thomas precession, velocity of the quark and its acceleration, respectively.

The absence of the spin-orbit force in the chirally restored regime is a very welcome feature because it is a well-known empirical fact that the spin-orbit force is either vanishing or very small in the spectroscopy in the u, d sector. In addition, for the rotating string $\vec{\sigma}(i) \cdot \vec{R}(i) = 0$, $\vec{\sigma}(i) \cdot \vec{R}(j) = 0$. The relations above immediately imply that the possible tensor interactions of quarks related to the string dynamics should be absent, once the chiral symmetry is restored.

9. Conclusions

We have demonstrated that the chiral symmetry of QCD is crucially important to understand the physics of hadrons in the u, d (and possibly in the u, d, s) sector. The low-lying hadrons are mostly driven by the spontaneous breaking of chiral symmetry. This breaking determines the physics and effective degrees of freedom in the low-lying hadrons. However, this physics is relevant only to the low-lying hadrons. In the high-lying hadrons chiral symmetry is restored, which is referred to as effective chiral symmetry restoration or chiral symmetry restoration of the second kind. A direct manifestation of the latter phenomenon is a systematic appearance of the approximate chiral multiplets of the high-lying hadrons. The essence of the present phenomenon is that the quark condensates which break chiral symmetry in the vacuum state (and hence in the low-lying excitations over the vacuum) become simply irrelevant (unimportant) for the physics of the highly excited states. The physics here is as if there was no chiral symmetry breaking in the vacuum. The valence quarks decouple from the quark condensates and consequently the notion of the constituent quarks with dynamical mass induced by chiral symmetry breaking becomes inadequate in the highly excited hadrons. This happens because in the highly excited hadrons the effects of quantum fluctuations vanish and a semiclassical regime is manifest.

Hence physics of the high-lying hadrons is mostly physics of confinement acting between the light quarks. Their very small current mass strongly distinguishes this physics from the physics of the heavy quarkonium, where chiral symmetry is irrelevant and the string (flux tube)

can be approximated as a static potential acting between the slowly moving heavy quarks. In the light hadrons in contrast the valence quarks are ultrarelativistic and their fermion nature requires them to have a definite chirality. Hence the high-lying hadrons in the u, d sector open a door to the regime of dynamical strings with chiral quarks at the ends. Clearly a systematic experimental exploration of the high-lying hadrons is required which is an interesting and important task and which should be of highest priority at the existing and forthcoming accelerators.

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